

The Weak Parity-Violating Pion-Nucleon Coupling

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Abstract

We use QCD sum rules to obtain the weak parity-violating pion-nucleon coupling constant $f_{\pi NN}$. We find that $f_{\pi NN} \approx 2 \times 10^{-8}$, about an order of magnitude smaller than the “best estimates” based on quark models. This result follows from the cancellation between perturbative and nonperturbative QCD processes not found in quark models, but explicit in the QCD sum rule method. Our result is consistent with the experimental upper limit found from ^{18}F parity-violating measurements.

In this Letter, we use the method of QCD sum rules with the electroweak and QCD Lagrangians to predict the weak parity-violating (PV) pion-nucleon coupling constant, $f_{\pi NN}$. The theoretical prediction of $f_{\pi NN}$ is an important and challenging problem. To-date, the most accurate PV experiments have only shown^{1,2)} that the upper limit for the magnitude of this coupling constant is 3-5 times smaller than the “best value” predicted by DDH³⁾ on the basis of a quark model and somewhat smaller than that in a similar calculation carried out more recently.⁴⁾ Since that time others have tried to estimate $f_{\pi NN}$ by means of chiral soliton models^{5,6)} and QCD sum rules.⁷⁾ This coupling is of particular interest because of its sensitivity to the neutral current contribution of weak nonleptonic processes at low energies.²⁾

QCD sum rules have been shown to be able to reproduce known properties of the nucleon, e.g., μ_p , μ_n , g_A , and of other hadrons.⁸⁾ However, they have rarely (if ever) been used to predict unknown properties. Keeping terms in the operator product expansion (OPE) up to dimension 5, we show that there are two main terms in the sum rule for $f_{\pi NN}$: the unit operator and a dimension D=3 susceptibility. By using an analogous sum rule for the strong coupling constant, $g_{\pi NN}$, to evaluate this susceptibility, we are able to determine the weak coupling $f_{\pi NN}$. An important aspect of the present work is that we demonstrate that there is a cancellation between perturbative and nonperturbative QCD modifications of the weak process.

We employ a two point function for the nucleon in an external pionic field. Our current is the usual one⁹⁾

$$\begin{aligned}\eta_p(x) &= \epsilon^{abc}[u^{aT}(x)C\gamma_\mu u^b(x)]\gamma^5\gamma^\mu d^c(x), \\ \bar{\eta}_p(y) &= \epsilon^{abc}[\bar{u}^b(y)\gamma_\nu C\bar{u}^{aT}(y)]\bar{d}^c(y)\gamma^\nu\gamma^5,\end{aligned}\tag{1}$$

where ϵ^{abc} is the antisymmetric tensor, C is the charge conjugation operator, and a, b, c are color indices. The neutron currents are similar, with the interchange of $d \leftrightarrow u$.

Since the most general weak PV π - N coupling is^{2,3,10)}

$$H_{PV}(\pi NN) = \frac{f_{\pi NN}}{\sqrt{2}}\bar{\psi}(\tau \times \phi_\pi)_3\psi,\tag{2}$$

only charged pions can be emitted or absorbed. For definiteness, we consider the absorption of a π^+ so that an initial neutron is converted to a proton, and the correlator we consider is

$$\Pi = i \int d^4x e^{ix \cdot p} \langle 0 | T[\eta_p(x) \bar{\eta}_n(0)] | 0 \rangle_{\pi^+} . \quad (3)$$

The general form of Π for the parity-violating pion-nucleon coupling, as dictated by relativistic invariance, is

$$\Pi^{PV} = \Pi_e 1 + \Pi_o \hat{p} , \quad (4)$$

with $\hat{p} \equiv \gamma_\mu p^\mu$.

The phenomenological evaluation of the correlator is carried out by deriving a dispersion relation for Π through the insertion of a complete set of physical intermediate states of spin $\frac{1}{2}$ in the expression of Eq. 3. Using the usual terminology, we refer to this as the right-hand side (RHS). We only use the sum rule for Π_e , since the sum rule for Π_o is not as stable. One finds for the parity-violating part of Eqs. 3,4:

$$\Pi_e^{PV}(p^2)^{RHS} = \frac{\lambda_N^2 f_{\pi NN} (p^2 + M^2)}{(p^2 - M^2)^2} + \textit{continuum}. \quad (5)$$

M is the nucleon mass; and the parameter λ_N is related to the amplitude for finding three quarks in a nucleon at one point and has been determined in a number of sum-rule calculations.⁸⁾ The double pole term, corresponding to the insertion of the one-nucleon intermediate state in Eq. (3), has contributions both from the weak pion-nucleon vertex and the parity violation in the nucleon state itself. As will be shown below, in our microscopic calculation using the two-point form only Z_0 -quark loops in the nucleon correlator give the parity-violating vertex correction. As is usual in the method, the physical property of interest, $f_{\pi NN}$, is obtained by treating the double-pole term explicitly, while the continuum and excited states are included in the numerical analysis via a parameterization, as discussed below.

The microscopic evaluation of Π , based on QCD and electroweak theory (the so-called left-hand sides (LHS) of the QCD sum rules for Π), is obtained by means of a Wilson coefficient expansion in inverse powers of p^2 . In this work we keep diagrams up

to dimension $D = 5$. The lowest dimensional diagrams which we consider are shown in Fig. 1. The higher dimensional diagrams which we include are obtained from those shown in Fig. 1 by the substitution of Figs. 2b and c for the pion-quark vertex, Fig. 2a. The propagators in coordinate space corresponding to the three diagrams of Fig. 2 are:

$$\begin{aligned} S_{5a}^{ab} &= \frac{i\vec{\tau} \cdot \vec{\pi}}{4\pi^2 x^2} g_{\pi q} \gamma^5 \delta^{ab} \\ S_{5b}^{ab} &= -\frac{i}{24} \vec{\tau} \cdot \vec{\pi} g_{\pi q} \chi_\pi < \bar{q}q > \delta^{ab} \gamma^5, \\ S_{5c}^{ab} &= \frac{i}{3 \cdot 2^7} m_0^\pi < \bar{q}q > g_{\pi q} \vec{\tau} \cdot \vec{\pi} x^2 \gamma^5, \end{aligned} \quad (6)$$

with $\chi_\pi g_{\pi q} \pi_j < \bar{q}q > \equiv < \bar{q} i \tau_j \gamma_5 q >_\pi$ and $m_0^\pi < \bar{q}q > \pi_j \equiv < \bar{q} i \gamma_5 g_c \tau_j \sigma \cdot G q >_\pi$. Here $g_{\pi q}$ is the pion-quark coupling, which is not explicitly used in the present calculation, and G represents the gluon field. The susceptibility χ_π enters in the evaluation of both strong and weak pion-nucleon coupling constants, while m_0^π enters only for the weak one. We will discuss the treatment of these parameters below. We only consider the even sum rule, namely that for Π_e ; that for Π_o involves further unknown susceptibilities. The evaluation of the diagrams is straightforward.

For the weak Hamiltonian, we take $H_w = \frac{G_F}{\sqrt{2}} (J^\mu J_\mu^\dagger + N^\mu N_\mu^\dagger)$ with

$$\begin{aligned} J^\mu &= \bar{u} \gamma^\mu (1 - \gamma_5) d \cos \theta_C \\ N^\mu &= \bar{u} \gamma^\mu (A_u + B_u \gamma^5) u + \bar{d} \gamma^\mu (A_d + B_d \gamma^5) d, \end{aligned} \quad (7)$$

where θ_C is the Cabibbo angle and A_u, A_d, B_u, B_d , are given by

$$\begin{aligned} A_u &= \frac{1}{2} \left(1 - \frac{8}{3} \sin^2 \theta_W\right), \\ A_d &= -\frac{1}{2} \left(1 - \frac{4}{3} \sin^2 \theta_W\right), \\ B_u &= -B_d = -\frac{1}{2}, \end{aligned} \quad (8)$$

with θ_W the Weinberg angle. This is the standard model Hamiltonian, which we use for the main part of the calculation. We then discuss the QCD effects on our results.

Since momentum can be transferred in the weak point-like interaction, shown by wavy lines representing Z^0 in the figures, there is an additional integral to be carried out in

the evaluation of Π . For example, we obtain for Fig. 1a

$$\begin{aligned}
\Pi_e^{1a} = & -2^6 G_F \sin^2 \theta_W g_{\pi q} \int d^D k_1 d^D k_2 d^D k_3 [k_1 \cdot (p - k_2 - k_3)(\hat{p} - \hat{k}_1 - \hat{k}_3)\hat{k}_2 \\
& + \frac{\epsilon}{4} \{2(k_2 \cdot (p - k_3)\hat{k}_1(\hat{p} - \hat{k}_3) + 2(p - k_1 - k_3) \cdot (p - k_2 - k_3)\hat{k}_2\hat{k}_1 \\
& + -3(\hat{p} - \hat{k}_1 - \hat{k}_3)\hat{k}_2\hat{k}_1(\hat{p} - \hat{k}_2 - \hat{k}_3)\}] \\
& [(2\pi)^{3D} k_1^2 k_2^2 k_3^2 (p - k_1 - k_2)^2 (p - k_2 - k_3)^2]^{-1}, \tag{9}
\end{aligned}$$

where $D = 4 - \epsilon$ is the dimension. There is no PV contribution from Figs. (1c) and (1d), and the sum of Figs. (1b) and (1e) vanish. The integrals in Eq. (9) are evaluated by standard Feynman techniques, with dimensional regularization. The result is

$$\Pi_e^{1a}(p^2) = -\frac{G_F \sin^2 \theta_W g_{\pi q}}{3^2 2^7 \pi^6} p^6 \ln(-p^2) \left(\frac{1}{\epsilon} + \frac{15}{2} - \frac{3}{2} \gamma \right). \tag{10}$$

We regularize the diagram using mass, vertex, and pion-quark vertex counter terms, leading to the one-loop corrections to our diagram shown in Fig. 3. The lowest dimension pion-quark vertex and mass renormalization diagrams for $f_{\pi NN}$ are shown in Figs. 3a-c. In our approximation of a contact weak interaction, the contribution of Figs. 3a-c vanish under a Borel transformation. The mechanism of Figs. 3d and 3e do not appear in the external field method. The only nonvanishing diagrams in the infinite Z-mass limit are those shown in Figs. 3f and 3g. With a minimal subtraction scheme we obtain an additional composite current, which we call η_V :

$$\begin{aligned}
\eta^V(p) &= \epsilon^{abc} [u^{aT}(k_1) C \gamma_\mu u^b(k_2)] \gamma^5 \Gamma_V^\mu d^c(k_3), \\
\Gamma_V^\mu &= \frac{4 G_F \sin^2(\theta_W)}{3^2 (4\pi)^2} (q^2)^{-\epsilon/2} (\hat{q} q^\mu - q^2 \gamma^\mu), \tag{11}
\end{aligned}$$

with $k_1 = p - k_2 - k_3$ and $q = k_2 + k_3$. This current is used for the vertex regularization shown in Figs. 3f and 3g. These vertex corrections give the contribution

$$\Pi_{e(V)}^{1a}(p^2) = \frac{G_F \sin^2 \theta_W g_{\pi q}}{3^2 2^7 \pi^6} p^6 \ln(-p^2) \left(\frac{1}{\epsilon} + \frac{14}{3} - \gamma \right). \tag{12}$$

Combining Eqs.(10,12) and taking the Borel transform one obtains for the regularized diagram 1a

$$\Pi_{e(R)}^{1a}(p^2) = \frac{G_F \sin^2 \theta_W g_{\pi q}}{3^2 2^7 \pi^6} \left(\frac{17}{3} - \gamma \right) M_B^8. \tag{13}$$

where M_B is the Borel mass. The other diagrams can be evaluated in the same manner. The results from the processes of Figs. 1 and those obtained from Figs. 1 with the substitution of Figs. 2b and c for Fig. 2a, with the counter terms given by corresponding substitution in Fig 3, are

$$\begin{aligned}
\Pi_e^{PV}(M_B^2) &= \frac{G_F \sin^2 \theta_W (\frac{17}{3} - \gamma)}{24\pi^2} M_B^4 [M_B^4 L^{-4/9} E_3 \\
&+ \frac{2}{3} \chi_\pi a L^{-4/9} M_B^2 E_2 + \frac{1}{2} m_0^\pi a E_1 L^{-4/9}] \\
&= f_{\pi NN} \bar{\lambda}_N^2 e^{-M^2/M_B^2} (2 \frac{M^2}{M_B^2} - 1)
\end{aligned} \tag{14}$$

where $a = -(2\pi)^2 < \bar{q}q >$ and $\bar{\lambda}_N^2 = (2\pi)^4 \lambda_N^2 / g_{\pi q}$. We do not include gluon condensate diagrams for $f_{\pi NN}$; they are of the same order or smaller than uncertainties of our calculation. The factors containing L , $L = 0.621 \ln(10M_B)$, give the evolution in Q^2 arising from the anomalous dimensions, and the $E_i(M_B^2)$ functions take into account excited states to ensure the proper large- M_B^2 behavior. The last line in Eq. (14) is the Borel transform of the double-pole term from the phenomenological (right-hand) side, Eq. (5). The direct proportionality to $\sin^2 \theta_W$ should be noted.

Finally, by explicit calculation or Fierz reordering, we can show that the contribution for W^\pm exchanges vanish. Thus, as required by symmetries^{3,10)}, we find no charged current contribution to the weak PV pion-nucleon vertex; such a contribution requires strangeness-changing currents and would thus be reduced by $\sin^2 \theta_C \approx 0.05$. Since we neglect strangeness in the nucleon and strangeness-changing currents, we obtain no contribution.

As we shall demonstrate below, the first two terms in the theoretical form for Π_e given in Eq. (14) are of opposite sign and tend to cancel. This is a crucial point. For this reason it is essential to either determine the value of the susceptibility χ_π from $g_{\pi NN}$ or to eliminate it from our equations. We do both as an aid in determining the stability of our solutions. First, we determine χ_π directly in terms of $g_{\pi NN}$ [as a function of the Borel mass] by using the sum rule for the strong coupling, which is analogous to Eq. (14), and attempt to use the result to determine $f_{\pi NN}$. Second, we eliminate χ_π from

the PV and strong coupling sum rules and find that we can determine $f_{\pi NN}$ in terms of $g_{\pi NN}$. Details are given below.

We use the correlator given by the two-point function of Eq. (3) for the strong as well as the weak interaction. The general form differs from Eq.(4) by the presence of a γ^5 in each term. The phenomenological (RHS) for the strong pion-nucleon coupling is now given by

$$\Pi_c^s(p^2)^{RHS} = \frac{\lambda_N^2 g_{\pi NN} M^2}{(p^2 - M^2)} \gamma^5 + \text{continuum}. \quad (15)$$

Unlike the weak PV pion-nucleon coupling, the evaluation of the strong one leads to a problem in that there is no double pole on the right-hand (dispersion relation) side. However, as shown by Reinders et. al.⁽¹¹⁾, the value of the coupling constant $g_{\pi NN}$ found in this way is virtually the same as that found by means of a 3-point function, which circumvents the lack of a double pole problem.

Keeping terms up to D=6, shown in Fig. 4, for the theoretical side (LHS), and taking the Borel transform we obtain the sum rule for the strong pion-nucleon coupling:

$$\begin{aligned} g_{\pi NN} \bar{\lambda}_N^2 e^{-M^2/M_B^2} &= M_B^6 L^{-4/9} E_2 - M_B^4 \chi_\pi a L^{2/9} E_1 + \frac{4}{3} a^2 L^{4/9} \\ &+ \frac{\langle g_c^2 G^2 \rangle E_0 M_B^2}{8} - \langle g_c^2 G^2 \rangle E_0 M_B^2 \left(\frac{13}{8} - \ln M^2 \right), \end{aligned} \quad (16)$$

where $\langle g_c^2 G^2 \rangle$ is the gluonic condensate.

Before we discuss our detailed evaluation of the sum rules to obtain our estimate of $f_{\pi NN}$, let us discuss the structure of Eqs. (14) and (16). First, as we discuss below, if we use the method of Ref. (12) [which uses arguments of PCAC within the sum rule context] to evaluate χ_π we find that $\chi_\pi a = -88 \text{ GeV}^2$. With this value, the χ_π term dominates both Eqs. (14,16) with the result that $g_{\pi NN} \simeq 155$ [in contrast to the experimental value of 13.5]. With this value of χ_π we find that $f_{\pi NN} \geq 10^{-6}$, at least an order of magnitude larger than experiment.

Secondly, since χ_π is the only unknown in Eq. (16), we can estimate the vacuum susceptibility using the experimental value of $g_{\pi NN} = 13.5$: this gives $\chi_\pi a \simeq -1.88 \text{ GeV}^2$,

two orders of magnitude smaller than the value given by the method of Ref. (12) [see discussion below]. With this value one finds that the first two terms in Eq. (14), the leading terms for $f_{\pi NN}$, almost cancel. Note that the second term involving χ_π enters with the opposite sign in the two equations for $f_{\pi NN}$ and $g_{\pi NN}$, respectively. This is the source of the very small parity-violating pion-nucleon coupling in comparison with quark model: there is a cancellation between the dimension zero model-like term using perturbative quark propagators and the vacuum pion susceptibility term.

However, we find that the sum rule obtained for $f_{\pi NN}$ [Eq.(14)], using the value of $\chi_\pi(M_B^2)$ extracted from Eq. (16), is not stable in M_B . Therefore we cannot obtain a reliable estimate of $f_{\pi NN}$ by this method.

We find that we can obtain a satisfactory sum rule to determine $f_{\pi NN}$ by eliminating χ_π from both Eq. (14) and Eq. (16) by taking derivatives with respect to M_B^2 . With this procedure, and taking the ratio of the weak to the strong sum rule we obtain the new sum rule for the weak in terms of the strong coupling constant:

$$\frac{f_{\pi NN}}{g_{\pi NN}} = c_w M_N^2 \frac{(M_N^2 - 4M_B^2)(E_3 M_B^4 + \frac{1}{2} a m_0^\pi E_1)}{(2M_N^4 + 3M_B^4 - 9M_N^2 M_B^2)(12E_2 M_B^4 + 3 \langle G^2 \rangle E_0)}, \quad (17)$$

where $c_w = G_F \sin^2 \theta_W (\frac{17}{3} - \gamma)/(24\pi^2) = 5.5 \times 10^{-8} \text{ GeV}^{-2}$. The sum rule is quite stable with a plateau in M_B^2 in the region expected, as shown in Fig. 5. Because of the strong cancellation between the first two terms in Eq. (14) [dimension 0 and dimension 2 terms], the dimension four term with the unknown parameter m_0^π is important for the final numerical value of $f_{\pi NN}$. We have taken $m_0^\pi = 0$ in Fig 5. Guided by the value of the parameter m_0 needed in the nucleon sum rule⁸⁾, we evaluate the sum rule given in Eq. 15 with m_0^π taken over the range 0.0 to +0.8 GeV. From this procedure we find:

$$\begin{aligned} f_{\pi NN} &\approx (1.9 \text{ to } 2.4) \times 10^{-8} \text{ for} \\ m_0^\pi &= (0 \text{ to } 0.8) \text{ GeV.} \end{aligned} \quad (18)$$

For negative values of m_0^π the value of $f_{\pi NN}$ becomes smaller and even negative, but we did not find stable solutions for sizable negative values of this unknown parameter. To

be consistent with the neglect of gluon condensate terms we quote as our central value of $f_{\pi NN}$ that with $m_0^\pi = 0$, shown in Fig. 5, as

$$f_{\pi NN} \approx 1.9 \times 10^{-8}. \quad (19)$$

This coupling constant is an order of magnitude smaller than the “best values” of Refs. 3 and 4. As emphasized earlier, this result follows from the cancellation of the two leading terms in Eq. (14). The first LHS term in that equation, a unit dimension term which would correspond to a quark model type calculation, gives a value for $f_{\pi NN} \approx 2 \times 10^{-7}$, similar to the quark model value. The second term, involving the nonperturbative QCD vacuum susceptibility, χ_π , strongly cancels the first term. Because of this cancellation, we cannot expect Eq. (19) to be very accurate, but we find a clear explanation for the small value of $f_{\pi NN}$, consistent with experiment.^{1,2)}

The results given in Eqs. 18 and 19 have been obtained using the Hamiltonian of the standard model (see Eqs. 7 and 8). Let us now consider the strong interaction modifications. These have been estimated in Refs. 3 and 4 using the renormalization group method. In the notation of Ref. 3, the operators involved in our calculation are O_4 and O_5 . Using the tables in Refs. 3 and 4 we find that the our parameter for the parity violation, $A_d B_u - A_u B_d$, would be changed by less than a factor of two in magnitude. Since the same parameter appears in all terms, this gives the overall uncertainty arising from strong interaction modifications. Therefore, the main conclusion of our work is not changed.

There are two relevant features that we would like to point out. The first one is that the use of pseudovector coupling also circumvents the problem of a lack of double pole for the strong interaction constant. For the Lagrangian

$$\mathcal{L}_{\pi NN} = \frac{g'_{\pi NN}}{m_\pi} \bar{\psi}_N i\gamma^\mu \gamma^5 \vec{\tau} \cdot \psi_N \nabla_\mu \vec{\phi}_\pi \quad (20)$$

we can treat $\nabla_\mu \phi_\pi$ as a constant external axial vector field. The QCD sum rule is then identical to our calculation of g_A .⁸⁾ At the quark level, we have

$$\mathcal{L}_{\pi qq} = \frac{1}{2f_\pi} \bar{\psi}_q i\gamma^\mu \gamma^5 \vec{\tau} \psi_q \nabla_\mu \vec{\phi}_\pi, \quad (21)$$

where f_π is the pion decay constant. From our previous result for g_A ,¹⁾ we then obtain

$$\begin{aligned}\frac{g'_{\pi NN}}{m_\pi} &= \frac{g_A}{2f_\pi}, \\ g_{\pi NN} &= g'_{\pi NN} \frac{2M}{m_\pi} = \frac{g_A M}{f_\pi}\end{aligned}\tag{22}$$

which is just the Goldberger-Treiman relation.

As a second feature we wish to attempt an independent estimate of χ_π . For this purpose we first use PCAC to obtain

$$\begin{aligned}\langle 0|\bar{u} i\gamma_5 u - \bar{d} i\gamma_5 d|\pi^0 \rangle &= \frac{-f_\pi m_\pi^2}{\sqrt{2}m_q} e^{-iq \cdot x} \\ &\equiv \bar{\chi}_\pi \phi_\pi < \bar{q}q > e^{-iq \cdot x},\end{aligned}\tag{23}$$

where we take the π -quark coupling to be unity in this discussion. We then use the work of Belyaev and Kogan¹²⁾, which assumes saturation of a sum by one pion states:

$$\begin{aligned}\langle 0|\bar{q} i\gamma_5 \tau_3 q|0 \rangle_\pi &= \frac{-i}{\sqrt{2}} \phi_\pi \int d^4x e^{iQ \cdot x} \langle 0|\bar{u} i\gamma_5 u - \bar{d} i\gamma_5 d|\pi \rangle < \pi|\bar{q} i\gamma_5 \tau_3 q|0 \rangle_{Q \rightarrow 0} \\ &= \frac{i}{\sqrt{2}} \phi_\pi \frac{f_\pi^2 m_\pi^2}{2m_q^2} \equiv \chi_\pi \phi_\pi < \bar{q}q >\end{aligned}\tag{24}$$

As described above, the value of χ_π obtained in this manner is more than an order of magnitude larger than that found by using the value of $g_{\pi NN}$ from experiment. Once more we point out that if we use it in Eq. (16) we find an order of magnitude discrepancy with the strong coupling constant, $g_{\pi NN}$. Furthermore, it is clear that this value of χ_π is inconsistent with Eq. (14), since by eliminating it with derivatives with respect to the Borel mass we obtain results an order of magnitude different than with its use. We conclude that Eq. (24) cannot be correct. We are not certain where the method of Belyaev and Kogan errs, but we believe that it is suspect. Note that $\chi_\pi = (f_\pi/m_q)\bar{\chi}_\pi \sim 20\bar{\chi}_\pi$

In conclusion, we find that the weak PV pion-nucleon coupling due to neutral currents is as small as that due to charged currents, $\sim 2 \times 10^{-8}$. This result agrees with the conclusion of the chiral soliton model of Kaiser and Meissner⁵⁾, but not that of Kaplan

and Savage⁶⁾. Our result also disagrees with quark model calculations^{3,4)} and with a previous QCD sum rule calculation.⁷⁾ If the coupling is as small as we estimate, it cannot be separated from the charged current contribution and thus cannot be found experimentally; and it is unlikely that the anapole will be seen.¹³⁾ Although we have omitted gluon condensate corrections to the PV correlator, our result is sufficiently small that these corrections will not alter our conclusion. Finally, we point out that in the two-point QCD sum rule method used here, the small value of $f_{\pi NN}$ which we obtained is the result of a cancellation between a process which can be treated in quark models and a vacuum process identified in the method of QCD sum rules.

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Figure Captions

Fig. 1. Lowest dimension quark diagrams for the PV weak pion-nucleon vertex. The dashed line represents a charged pion and the wavy line a Z^0 .

Fig. 2. Quark propagator modifications in an external pion field.

Fig. 3. Pion-nucleon weak vertex correction diagrams.

Fig. 4. Diagrams contributing to the calculation of $g_{\pi NN}$.

Fig. 5. Solution for $f_{\pi NN}$ from Eq. 15 as a function of M_B .

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